

Tight bonds between sterile neutrinos and dark matter

Torsten Bringmann,^a Jasper Hasenkamp^b and Jörn Kersten^c

^aDepartment of Physics, University of Oslo,
Box 1048, Oslo, NO-0316 Norway

^bCenter for Cosmology and Particle Physics, New York University,
70 Washington Square S, New York, NY, 10003 U.S.A.

^cII. Institute for Theoretical Physics, University of Hamburg,
Luruper Chaussee 149, Hamburg, DE-22761 Germany

E-mail: Torsten.Bringmann@fys.uio.no, Jasper.Hasenkamp@nyu.edu,
Joern.Kersten@desy.de

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Abstract. Despite the astonishing success of standard Λ CDM cosmology, there is mounting evidence for a tension with observations at small and intermediate scales. We introduce a simple model where both cold dark matter (DM) and sterile neutrinos are charged under a new $U(1)_X$ gauge interaction. The resulting DM self-interactions resolve the tension with the observed abundances and internal density structures of dwarf galaxies. At the same time, the sterile neutrinos can account for both the small hot DM component favored by cosmological observations and the neutrino anomalies found in short-baseline experiments.

Keywords: dark matter theory, cosmological neutrinos, cosmological parameters from LSS, cosmological parameters from CMBR

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1 Introduction

The nature of the dark matter (DM) in the Universe is one of the outstanding puzzles of science. Intensive research over the past decades led to the current standard DM paradigm of cold, collisionless particle DM (CDM). Alongside the remarkable success of the standard cosmological model, known as Λ CDM, in describing the Universe at large scales, a number of problems became increasingly solid over the past years driven by the increase in precision of cosmological and astrophysical observations.

On small cosmological scales those are most prominently: 1) *missing satellites* — the DM halo of the Milky Way (MW) should contain many more dwarf-sized subhalos (satellites) than observed [1, 2]. Beyond the MW, observed galaxy luminosity and HI-mass functions show shallower faint-end slopes than predicted [3]. 2) *Cusp vs. core* — low surface brightness and dwarf galaxies seem to have cored inner density profiles, at odds with CDM cusps predicted by simulations [4, 5]. 3) *Too big to fail* — the observed brightest satellites of the MW attain their maximum circular velocity at too large radii in comparison to the densest and most massive satellites found in simulations, i.e. the latter have no counterpart in observations even though they should be very efficient in forming stars [6, 7]. Both astrophysical [8–16] and DM-related [17–32] solutions are actively discussed, but the first *simultaneous* solution has only been advanced rather recently [33] (see also ref. [34] for a two-component DM model that augurs well in this respect).

On intermediate scales the observed galaxy cluster mass function and galaxy shear power spectra, both tracing the matter power spectrum, are inconsistent with observations of the cosmological microwave background (CMB). Also the current expansion rate of the Universe H_0 appears inconsistent with CMB data. More generally, “local” measurements of cosmological parameters, probing cosmology at “low” redshifts $z \lesssim 10$, are inconsistent with CMB observations that are particularly sensitive to high redshifts $z \gtrsim 1000$ [35]. A mixed DM model, which adds a small hot dark matter (HDM) component to the dominant CDM, resolves those inconsistencies and thus reconciles the low and high redshift universe [36–39].

This requires an effective number of additional neutrino species ΔN_{eff} and an effective HDM mass of [37]

$$\Delta N_{\text{eff}}|_{\text{cmb}} = 0.61 \pm 0.30, \quad m_{\text{hdm}}^{\text{eff}} = (0.41 \pm 0.13) \text{ eV}. \quad (1.1)$$

Here, we extend the standard model of particle physics (SM) by a spontaneously broken $U(1)_X$ gauge theory and introduce a sterile neutrino that is not charged under SM interactions. We demonstrate that all aforementioned problems of standard cosmology are resolved by coupling both CDM and sterile neutrinos, the latter automatically being promoted to the desired HDM component, to a $U(1)_X$ gauge boson of $\mathcal{O}(\text{MeV})$ mass. This can be achieved for parameters that resolve the anomalies [40–46] reported in short-baseline neutrino oscillation experiments, in particular a sterile neutrino mass of $\sim 1 \text{ eV}$.

This article is organized as follows. We first set up our model and give a brief overview of the cosmological evolution of the non-standard degrees of freedom introduced here, as sketched in figure 1. We continue by showing that the properties of the thermally produced CDM sector can provide a solution to all ΛCDM problems at small scales. Next, we demonstrate that our sterile neutrino HDM component can *simultaneously* satisfy the cosmologically favored values stated above and describe the anomalies in short-baseline oscillation experiments. We conclude with a discussion and an outlook.

2 Model setup

We consider the extension of the SM gauge group, $G_{\text{SM}} = \text{SU}(3)_c \times \text{SU}(2)_L \times \text{U}(1)_Y$, by an Abelian gauge symmetry $U(1)_X$ with corresponding gauge boson V . We introduce a Dirac fermion χ at the TeV scale, which will form the CDM, and two right-handed neutrinos $\nu_{R_{1,2}}$. Those new particles are neutral under G_{SM} but carry $U(1)_X$ charges, while the SM particles are neutral under $U(1)_X$. Anomaly cancellation requires the $\nu_{R_{1,2}}$ to carry charges of opposite sign with equal absolute value; for concreteness, we take the charges of $(\chi, \nu_{R_1}, \nu_{R_2})$ to be $(1, X_{\nu_R}, -X_{\nu_R})$.

We further assume that the $U(1)_X$ is spontaneously broken at the MeV scale by the vacuum expectation value (VEV) v_Θ of a complex Higgs field Θ , which is a representation $(1, 0, 2X_{\nu_R})$ under $\text{SU}(2)_L \times \text{U}(1)_Y \times \text{U}(1)_X$, while the Higgs field ϕ responsible for the electroweak symmetry breaking is a $(2, 1/2, 0)$. Another complex scalar ξ , with charges $(1, 0, X_{\nu_R})$ and VEV $v_\xi < v_\Theta$, is introduced to enable active-sterile neutrino mixing.

After symmetry breaking the low-energy, effective Lagrangian of our theory reads

$$\mathcal{L} = \mathcal{L}_{\text{SM}} + \mathcal{L}_R + \mathcal{L}_\chi + \mathcal{L}_{\text{kin. mix.}} + \mathcal{L}_{\text{Higgs}}. \quad (2.1)$$

Here, \mathcal{L}_{SM} denotes SM terms and \mathcal{L}_R contains

$$\mathcal{L}_R \supset -\frac{1}{2}\overline{\nu_{R_1}^c} M_1 \nu_{R_1} - \frac{1}{2}\overline{\nu_{R_2}^c} M_2 \nu_{R_2} - \overline{\nu_{R_1}^c} M_{RR} \nu_{R_2} - \overline{\nu_L} M_{LR} \nu_{R_1} + \text{h.c.}, \quad (2.2)$$

in addition to kinetic terms and Majorana mass terms for the SM neutrinos. The active-sterile neutrino mixing arises from a dimension-5 operator with $M_{LR} \sim v_\phi v_\xi / \Lambda$, suppressed by a scale Λ defined by the UV completion of the theory. The mass eigenstates $(\nu_1, \nu_2, \nu_3, N_1, N_2)$ are mixtures of the flavor eigenstates $(\nu_e, \nu_\mu, \nu_\tau, \nu_{R_1}^c, \nu_{R_2}^c)$. With the short-baseline anomalies in mind, Yukawa couplings between $\nu_{R_{1,2}}$ and Θ are chosen such that $m_{N_2} \sim M_2 \sim \text{MeV} \gg m_{N_1} \sim M_1 \sim \text{eV}$. We note that even a very small mixing between ν_{R_1} and ν_{R_2} , with $M_{RR}/M_2 \gtrsim 10^{-6}$, allows the cosmologically fast decay of ν_{R_2} .

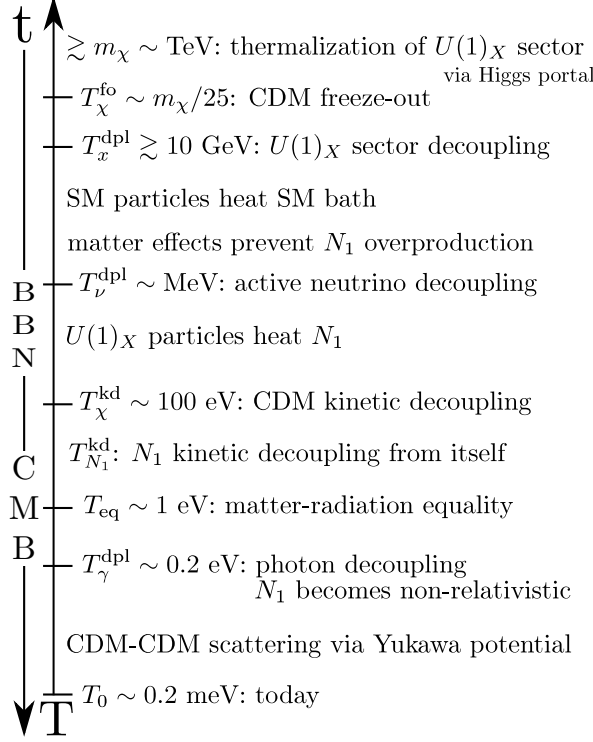


Figure 1. Schematic overview of the cosmology implied by the model defined in eq. (2.1).

Terms in \mathcal{L} related to V and the new fermions are

$$\mathcal{L}_x = \bar{\chi} (i\not{\partial} - m_\chi) \chi - \frac{1}{4} F_{\mu\nu}^x F^{x\mu\nu} - \frac{1}{2} m_V^2 V_\mu V^\mu - g_X V_\mu (X_{\nu_R} \bar{\nu}_{R1} \gamma^\mu \nu_{R1} - X_{\nu_R} \bar{\nu}_{R2} \gamma^\mu \nu_{R2} + \bar{\chi} \gamma^\mu \chi), \quad (2.3)$$

where g_X denotes the $U(1)_X$ gauge coupling. To ensure the stability of χ we might impose a discrete Z_2 symmetry under which only χ is assigned a negative parity. The symmetries also allow a kinetic mixing term $\mathcal{L}_{\text{kin. mix.}} = -\frac{\epsilon}{2} F_{\mu\nu}^x F^{\mu\nu}$, where $F^{x\mu\nu}$ ($F^{\mu\nu}$) denotes the $U(1)_X$ (electromagnetic) field strength tensor. We assume $\epsilon \ll 1$ to satisfy the severe existing constraints on this parameter [47, 48].

We refer to ref. [49] for a general discussion of the Higgs sector for Θ and ϕ as contained in $\mathcal{L}_{\text{Higgs}}$, adopting that m_V and the mass of the new light Higgs boson h_x are of the same order of magnitude in the relevant cases. The ‘‘Higgs portal’’ term in $\mathcal{L}_{\text{Higgs}} \supset \kappa |\phi|^2 |\Theta|^2 \supset \frac{\kappa}{4} v_\phi \phi \Theta^2 \simeq \frac{\kappa}{4} v_\phi h h_x^2$, where we have assumed a negligible mixing between h_x and the SM-like Higgs h in the last step, connects the SM and the new $U(1)_X$ sector. For simplicity, we assume the couplings of the additional portal terms $|\phi|^2 |\xi|^2$ and $|\Theta|^2 |\xi|^2$ to be negligibly small.

3 Thermalization via the Higgs portal and decoupling of the dark sector

In figure 1 we provide a schematic overview of the cosmology arising from our model represented by eq. (2.1). Before electroweak symmetry breaking, the 4-scalar interaction $\kappa |\Theta|^2 |\phi|^2$ keeps the $U(1)_X$ sector in thermal equilibrium with the SM bath if the thermalization rate $\Gamma_{\text{th}} = n_{h_x} \langle \sigma_{\text{th}} v_{\text{rel}} \rangle \sim 10^{-3} \kappa^2 T$ is larger than the expansion rate $H \sim 10 T^2 / M_{\text{pl}}$. In those

expressions, n_{h_x} denotes the number density of h_x and $\langle \sigma_{\text{th}} v_{\text{rel}} \rangle$ the thermally averaged annihilation cross section of h_x pairs. If we, e.g., require thermal equilibrium at temperatures below 10 TeV, i.e. above the CDM mass m_χ , we obtain a lower bound on the Higgs portal coupling of $\kappa \gtrsim 10^{-6}$. After electroweak symmetry breaking the relevant process becomes $h_x h_x \xrightarrow{h} f \bar{f}$, controlled by the $\frac{\kappa}{4} v_\phi h h_x^2$ coupling. The thermalization rate is then $\Gamma_{\text{th}} \sim 10^{-3} \kappa^2 T^3 m_f^2 / m_h^4$, with f corresponding to the heaviest relativistic SM fermion, so the decoupling temperature becomes $T_x^{\text{dpl}} \sim 10^3 m_h^4 / (\kappa^2 m_f^2 M_{\text{pl}})$. For details on thermalization via the Higgs portal we refer to [50–52], where thorough calculations of h_x abundances for $m_{h_x} \sim \text{TeV}$ were performed (while in our case h_x decouples relativistically).

The particles in the dark sector are tightly coupled to each other due to the $U(1)_X$ interaction, and more weakly to the SM via the Higgs portal. Once the latter ceases to be effective, the whole $U(1)_X$ sector therefore decouples from the SM bath and entropy is conserved separately in the two sectors. Whenever a particle in equilibrium becomes non-relativistic it thus heats its bath, thereby increasing the temperature by a factor $\left(g_{*,\nu/N_1}^{\text{before}} / g_{*,\nu/N_1}^{\text{after}} \right)^{\frac{1}{3}}$, where $g_{*,i}$ counts the effective degrees of freedom (d.o.f.) determining the entropy density of the sector in thermal equilibrium with the species i . The non-standard contribution to the radiation density is then given by

$$\Delta N_{\text{eff}}(T) = \frac{T_{N_1}^4}{T_\nu^4} = \left(\frac{g_{*,\nu}}{g_{*,N_1}} \right)^{\frac{4}{3}} \bigg|_T \left(\frac{g_{*,N_1}}{g_{*,\nu}} \right)^{\frac{4}{3}} \bigg|_{T_x^{\text{dpl}}} . \quad (3.1)$$

The *maximal* possible value of this quantity at the onset of big bang nucleosynthesis (BBN), at $T \sim 1 \text{ MeV}$, is then obtained if *all* new particles but the light sterile neutrino, N_1 , have become non-relativistic by then. This results in

$$\Delta N_{\text{eff}}^{\text{max}}|_{\text{bbn}} \simeq \left[58.4 / g_{*,\nu} \left(T_x^{\text{dpl}} \right) \right]^{\frac{4}{3}} , \quad (3.2)$$

well within bounds from BBN [53–55] for $T_x^{\text{dpl}} \gtrsim 1 \text{ GeV}$.

4 Self-interacting CDM

At high temperatures, the DM particles are kept in chemical equilibrium via $\chi\chi \leftrightarrow VV$ (for unit sterile neutrino charges, $X_{\nu_R} \sim 1$, also the annihilation into $\bar{\nu}_R \nu_R$, $h_x h_x$ and $\xi^* \xi$ via a virtual V becomes important). For TeV-scale DM the number density freezes out at sufficiently early times ($T_\chi^{\text{fo}} \sim m_\chi / 25$) to still have $T_V = T$. Assuming for simplicity $X_{\nu_R} \ll 1$, the CDM relic density then becomes

$$\Omega_{\text{cdm}} h^2 = 2 \Omega_\chi h^2 \sim 0.11 \left(\frac{0.67}{g_X} \right)^4 \left(\frac{m_\chi}{\text{TeV}} \right)^2 \quad (4.1)$$

up to $\mathcal{O}(1)$ corrections due to the Sommerfeld effect [58], which we fully take into account [33]. This fixes g_X for a given m_χ throughout this work.

Kinetic decoupling [59] of χ happens much later and is determined by the elastic scattering rate for $\chi N_1 \leftrightarrow \chi N_1$. For a thermal distribution of sterile neutrinos, the decoupling temperature is given by [33]

$$T_\chi^{\text{kd}} \simeq \frac{62 \text{ eV}}{X_{\nu_R}^{\frac{1}{2}} g_X} \left(\frac{T}{T_{N_1}} \right)_{\text{kd}}^{\frac{3}{2}} \left(\frac{m_\chi}{\text{TeV}} \right)^{\frac{1}{4}} \left(\frac{m_V}{\text{MeV}} \right) , \quad (4.2)$$

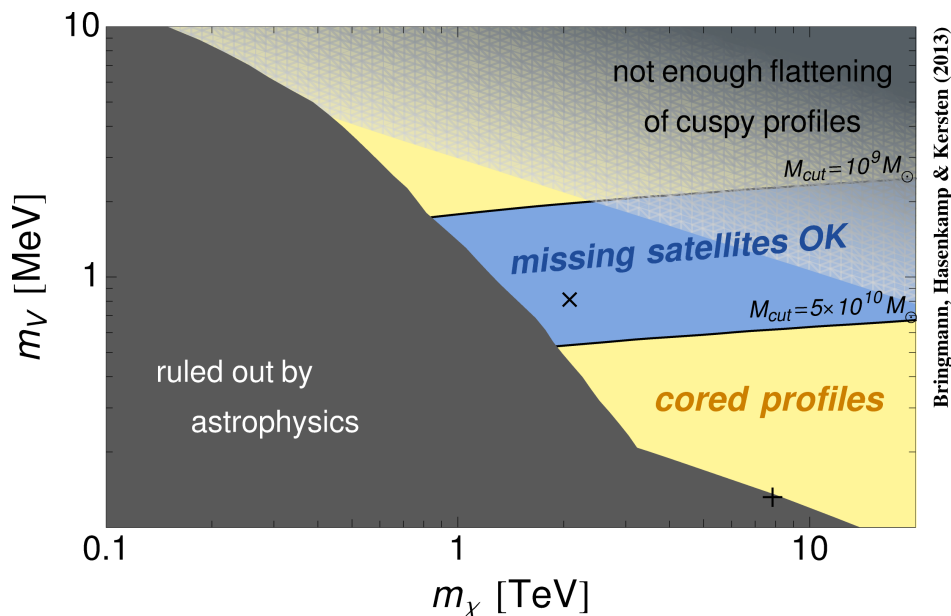


Figure 2. In the yellow area, the CDM self-interaction is strong enough to flatten density *cusps* in the inner parts of (dwarf) galaxies [30] and likely also solves the *too big to fail* problem (as explicitly demonstrated in N -body simulations for parameter values corresponding to the crosses [31]). The dark area is excluded by astrophysics [29, 30, 56, 57]. The blue band addresses the *missing satellites* problem [33], with a normalization that — according to eq. (4.2) — is proportional to $m_V \propto X_{\nu_R}^{1/2} (T_{N_1}/T)_{\text{kd}}^{3/2}$. Here, we show for reference the case of $X_{\nu_R} = 0.2$ and $(T_{N_1}/T)_{\text{kd}} = 0.46$.

which translates into a cutoff in the power spectrum of matter density perturbations at $M_{\text{cut}} \sim 1.7 \times 10^8 (T_{\chi}^{\text{kd}}/\text{keV})^{-3} M_{\odot}$. We note that the light mass eigenstates ν_i also acquire a $U(1)_X$ charge from their ν_R^c component; this will further lower T_{χ}^{kd} if $\sin \theta \gtrsim (T_{N_1}/T_{\nu})^{3/2}$.

After structure formation, the $U(1)_X$ -induced Yukawa potential produces galaxy cores that match the observed velocity profiles of massive MW satellites, solving *cusp vs. core* [28, 30] and *too big to fail* [31], while avoiding constraints on DM self-interactions on larger scales [30]. At the same time, the late kinetic decoupling addresses the *missing satellites* by suppressing the matter power spectrum at dwarf galaxy scales [33] (see also refs. [24, 60–62]). In figure 2, we show the desired parameter space for m_V and m_{χ} (based on ref. [33], but using an improved parameterization [63] of the Yukawa scattering cross section [28, 64–67]). The blue band, in particular, shows the range of masses that allow a solution of the missing satellites problem. Note that its normalization depends on the choice of X_{ν_R} and $(T_{N_1}/T)_{\text{kd}}$, whereas the m_{χ} -dependence is uniquely determined by eq. (4.2) and the form of $g_X(m_{\chi})$ that corresponds to the correct thermal CDM relic density, cf. eq. (4.1). The dark area is excluded by the requirements to not disrupt galactic satellites and to avoid a gravothermal catastrophe [29, 30, 56, 57].

5 The HDM component

We will now address the question whether the N_1 population in our model can account for the cosmologically preferred HDM component [36–39]. In the absence of any significant

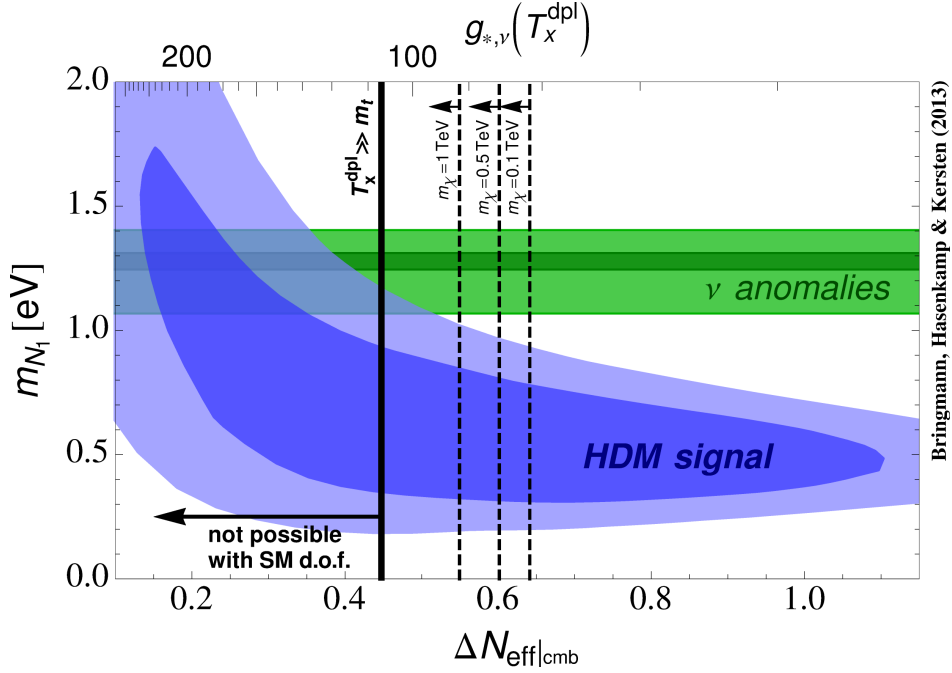


Figure 3. Sterile neutrino mass m_{N_1} vs. late-time additional relativistic d.o.f. $\Delta N_{\text{eff}}|_{\text{cmb}}$ and SM d.o.f. at decoupling of the $U(1)_X$ sector, cf. eq. (3.2). Shaded areas correspond, at 1σ and 2σ respectively, to the HDM signal [37] and values of m_{N_1} favored by the neutrino anomalies [68, 69]. Dashed lines indicate the minimal value of $\Delta N_{\text{eff}}|_{\text{cmb}}$ compatible with a CDM mass of, from right to left, $m_\chi = 100, 500, 1000$ GeV. Parameter values to the left of the solid line are not achievable in the minimal scenario studied here.

additional N_1 production mechanism, see the discussion further down, we simply have

$$\Delta N_{\text{eff}}|_{\text{cmb}} = \Delta N_{\text{eff}}^{\text{max}}|_{\text{bbn}}. \quad (5.1)$$

From the definition $m_{\text{hdm}}^{\text{eff}} \equiv [T_{N_1}/T_\nu^{\text{ACDM}}]^3 m_{N_1} = 11/4 [T_{N_1}/T]^3 m_{N_1}$ it furthermore follows that

$$m_{\text{hdm}}^{\text{eff}} = (\Delta N_{\text{eff}}|_{\text{cmb}})^{3/4} m_{N_1}. \quad (5.2)$$

By choosing the right decoupling temperature in eq. (3.2), which in our model corresponds to adjusting κ , we can then reproduce eq. (1.1). This is demonstrated in figure 3 where we show the allowed region of $\Delta N_{\text{eff}}|_{\text{cmb}}$ and $m_{\text{hdm}}^{\text{eff}}$ [37] in terms of $g_{*,\nu}(T_x^{\text{dpl}})$ and m_{N_1} .

The thermal production of the CDM component as treated here requires $T_x^{\text{dpl}} \lesssim T_\chi^{\text{fo}} \sim m_\chi/25$. On the other hand, $g_{*,\nu}(T_x^{\text{dpl}})$ cannot exceed the full number of SM d.o.f. even for very early $U(1)_X$ decoupling. Taken together, this points to $0.2 \text{ eV} \lesssim m_{N_1} \lesssim 1.2 \text{ eV}$.

6 Neutrino anomalies

Oscillation experiments observing neutrinos from accelerators [40, 41], reactors [42, 43] (but see [70]), and radioactive sources [44–46] reported anomalies that may indicate the existence of sterile neutrinos with a mass squared difference $\Delta m^2 \sim 1 \text{ eV}^2$ to the SM neutrinos. In figure 3, we show the 1σ and 2σ ranges for Δm^2 from [68, 69] for orientation, assuming

$m_{N_1}^2 = \Delta m^2$. These ranges were obtained from a global fit of oscillation data assuming the existence of a single sterile neutrino (note that it is being debated to what extent the data can be consistently explained by oscillations alone and whether a second sterile neutrino is necessary to achieve a satisfactory fit [71], which would not be possible to accommodate in our setup). From figure 3 we find that the regions allowed by the HDM signal and neutrino oscillations indeed overlap. While this happens only at the 2σ level, we note that the corresponding range of ΔN_{eff} is the *same* as independently suggested by $m_\chi \gtrsim 1$ TeV (as favored from figure 2).

7 Discussion

Before standard neutrino decoupling at $T \sim 1$ MeV, the effective mixing angle θ_m between active and sterile neutrinos is strongly suppressed due to the matter potential generated by the $U(1)_X$ couplings of the sterile neutrinos [72] (see also [73]). As the Universe cools, the effective mixing angle eventually reaches its vacuum value θ . This may give rise to an additional production of sterile neutrinos due to their $U(1)_X$ interaction. The largest effect on the scenario sketched above would result if the neutrinos completely re-thermalized, creating a thermal N - ν bath.

In that case, conservation of entropy density allows us to determine the temperature $T_{N\nu}$ of the newly established neutrino bath as $4T_{N\nu}^3 = \left[N_{\text{eff}}^{\text{SM}} + (\Delta N_{\text{eff}}|_{\text{bbn}}^{\text{max}})^{\frac{3}{4}} \right] (T_\nu^{\text{ACDM}})^3$, where $N_{\text{eff}}^{\text{SM}} \simeq 3.046$. Rather than eqs. (5.1), (5.2), we thus obtain

$$\Delta N_{\text{eff}}|_{\text{cmb}} = \frac{1}{4^{\frac{1}{3}}} \left[N_{\text{eff}}^{\text{SM}} + (\Delta N_{\text{eff}}|_{\text{bbn}}^{\text{max}})^{\frac{3}{4}} \right]^{\frac{4}{3}} - N_{\text{eff}}^{\text{SM}}, \quad (7.1)$$

$$m_{\text{hdm}}^{\text{eff}} = \frac{1}{4} \left[N_{\text{eff}}^{\text{SM}} + (\Delta N_{\text{eff}}|_{\text{bbn}}^{\text{max}})^{\frac{3}{4}} \right] m_{N_1}. \quad (7.2)$$

Rewriting this as $m_{N_1} = 2\sqrt{2} m_{\text{hdm}}^{\text{eff}} / (\Delta N_{\text{eff}}|_{\text{cmb}} + N_{\text{eff}}^{\text{SM}})$, we immediately see that in the re-thermalization case a sterile neutrino can still consistently explain the HDM signal — but only if its mass is considerably smaller than required by the neutrino anomalies.

Turning to potential constraints on our scenario, BBN limits are easily satisfied as already stressed earlier. CDM also decouples kinetically too early to imprint observable dark acoustic oscillation (DAO) features in the CMB [74]. Final state V radiation in the decay of SM particles [75], finally, does not constrain our scenario because V does not couple to left-handed neutrinos. An interesting aspect of our HDM component is that it does not necessarily manifest itself as perfectly free-streaming particles in the CMB or during structure formation, which in principle can be probed [76]; by comparing the elastic scattering rate with the Hubble expansion, we rather expect complete decoupling only at $T_{\nu_R\nu_R}^{\text{dpl}} \sim 3 \text{ eV} \left(\frac{m_\chi}{\text{TeV}} \right)^{-\frac{2}{3}} \left(\frac{X_{\nu_R}}{0.2} \right)^{-\frac{2}{3}} \left(\frac{m_V^2/X_{\nu_R}}{\text{MeV}^2/0.2} \right)^{\frac{2}{3}}$, where the last factor must be of order unity (see figure 2).

The dominant decay channel of our sterile neutrino is $N_1 \rightarrow \nu\nu\nu$. Even though this is strongly enhanced compared to the analogous common decay via a virtual Z [77], we find the resulting lifetime for the best-fit neutrino mixings [68] to be

$$\tau_{N_1} \sim 10^5 t_0 \left(\frac{m_\chi}{\text{TeV}} \right)^{-2} \left(\frac{X_{\nu_R}}{0.2} \right)^{-2} \left(\frac{m_V^2/X_{\nu_R}}{\text{MeV}^2/0.2} \right)^2 \left(\frac{\text{eV}}{m_{N_1}} \right)^5, \quad (7.3)$$

which greatly exceeds the age of the Universe t_0 . We note that the decay $N_1 \rightarrow \nu\gamma$ is even more suppressed due to the necessarily small value of ϵ .

8 Conclusions

In this article we have considered a mixed DM model as favored by recent cosmological observations, which adds a small HDM component to the dominant CDM, the former consisting of an eV-scale sterile neutrino and the latter of a TeV-scale Dirac fermion. We have studied the cosmological consequences of equipping both these particles with charges under a new spontaneously broken $U(1)_X$ gauge theory, under which all SM particles are singlets.

Thermalizing the $U(1)_X$ sector in the early universe via the so-called *Higgs portal* allows the thermal production of the CDM. The sterile neutrinos would also be thermally produced and elegantly form the HDM component, essentially because the $U(1)_X$ sector decouples much earlier than SM neutrinos. Remarkably, this is possible for a set of parameters that equip the CDM particles with a $U(1)_X$ mediated self-interaction that is of the right form and magnitude to provide a simultaneous solution to the small-scale problems of Λ CDM cosmology [33]. Finally, overproduction via mixings is likely prevented by the large thermal potential that the sterile neutrinos create by their $U(1)_X$ interactions [72, 73]; in this case one can even address the neutrino oscillation anomalies within the same framework. In other words, a sterile neutrino as preferred by neutrino oscillation anomalies would not only be reconciled with cosmology but promoted to the desired HDM component.

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Note added. After the arXiv submission of this work, several further concrete models have been proposed that simultaneously address all shortcomings of Λ CDM on small scales [72, 78, 79], based to a varying degree on the general mechanism suggested in ref. [33].

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